Direct mechanism in
solar nuclear reactions

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Abstract. A short overview of the direct reaction mechanism and the models used for the analysis of such processes is given. Nuclear reactions proceeding through the direct mechanism and involved in solar hydrogen burning are discussed. The significance of these nuclear reactions with respect to the solar neutrino problem is investigated.

INTRODUCTION

In the last years the importance of the direct reaction (DI) mechanism in stellar nuclear reactions has been realized. This mechanism is of relevance for many nuclear reactions in primordial and stellar nucleosynthesis. Obviously, the best stellar data are obtained from the observation of the sun. In this article we want to concentrate on direct nuclear reactions taking place in the pp–chain of solar hydrogen burning and being related to the solar neutrino problem; i.e. that the number of solar neutrinos observed by earth–bound detectors is significantly less than predicted by the Standard Solar Model (SSM).

This article concentrates more on solar reactions rates obtained from calculational efforts than from experimental determination. In section 2 we give a short overview of the DI mechanism and models. These models are *microscopic methods*, as the Resonating Group Method (RGM) and Generator Coordinate Method (GCM) and *potentials models*, as the Optical Model (OM), Distorted Wave Born Approximation (DWBA) and Direct Capture (DC). In section 3 we discuss some nuclear reactions which proceed predominately through the direct reaction mechanism and are of interest for the solar neutrino problem. In the following section we discuss some nuclear aspects of the solar neutrino problem. Finally, in the last section we give a summary.

DIRECT REACTION MECHANISMS AND MODELS

In nuclear reactions two extreme types of reaction mechanisms can exist: The compound–nucleus (CN) and direct (DI) process. In the CN mechanism the projectile merges in the target nucleus and excites many degrees of freedom of the CN. The excitation proceeds via a multistep process and therefore has a reaction time typically of the order 10^{-16} s to 10^{-20} s. After this time the CN decays into various exit channels. The relative importance of the decay channels is determined by the branching ratios to the final states. In the DI process the projectile excites only a few degrees of freedom (e.g. single–particle or collective). The excitation proceeds in one single step and has a characteristic time scale of 10^{-21} s to 10^{-22} s. This corresponds to the time the projectile needs to pass through the target nucleus; this time is much shorter than the reaction time of CN processes. However, at subCoulomb energies the reaction is now hindered by the Coulomb and centrifugal barriers. Therefore, the characteristic time scale is enhanced by a factor determined by the barrier penetration probabilities.

The question whether a given reaction favours CN or DI processes depends on the reaction considered and on the relative energies in the entrance and exit channel. In general, a given reaction involves both types of reaction mechanisms and also intermediate types (e.g. precompound reactions). However, for certain reactions and projectile energy ranges one type of reaction mechanism may dominate.

In thermonuclear scenarios the projectile energy is well below the Coulomb and/or centrifugal barrier. At these energies the competition between different reaction mechanisms is quite complicated. At these energies the CN formation may be suppressed, because there may exist no CN levels that can be populated, especially in light or magic nuclei. If the overlap between entrance– and exit–channel wave functions is large, then the DI process will be enhanced. This may be the case if the wave function of the transferred nucleon or nucleon cluster is weakly bound. Then the reactions may take place well outside the nucleus. On the other hand in this case the CN process is suppressed, because the overlap with the confined bound–state type wave functions which are responsible for the CN process, is small.

Closely connected to the DI and CN mechanisms are the resonance types in the nuclear reaction cross sections. There are again two extreme cases for resonance structures. The first type (compound resonance) occurs in CN reactions and has a narrow width (order of 1 to 100 eV). This width is determined by the long reaction time corresponding to the many steps involved in the formation and decay of the CN. The second type is called potential, optical, single–particle, potential-well or scattering resonance. This resonances type is associated with DI and is determined by the optical potential in the entrance or exit channel. Usually the widths of potential resonances are broad (order

of 10 keV to a few MeV) corresponding to the short reaction times of the DI. These resonances lose their meaning for higher energies because of their broad widths, as a result of absorption into other channels. However, for energies below the Coulomb and centrifugal barriers the width of potential resonances may become very narrow due to increasing lifetimes.

First–principle microscopic theories, such as the resonating group method (RGM) or generator coordinate method (GCM), are based on many–nucleon wave functions of the nuclei involved and on nucleon–nucleon interactions (1), (2). In this approach, the explicit inclusion of the Pauli principle leads to complicated highly nonlocal potentials for the interaction between the composite nuclei in the entrance and exit channel.

The main drawback of the RGM is that it requires extensive analytical calculations without systematic character when going from one reaction to another. Consequently, the application of the RGM is essentially restricted to reactions involving only a small number of nucleons. This problem can be overcome by the GCM which is similar to the RGM, but it allows systematic calculations, well adopted to a numerical approach. In the GCM the relative wave functions are expanded in a Gaussian basis. Nowadays, most microscopic calculations are performed in the GCM framework.

It is obvious that a fully microscopic approach like RGM and GCM is more satisfying, since it is a first–principle approach. Such a microscopic model starts from the nucleon–nucleon interaction and does not contain any free parameters. It is therefore possible to predict physical properties of the system independently of experimental data. However, in most cases the fully microscopic approach rarely reproduces physical quantities which are fundamental for the calculation of astrophysical cross sections, such as thresholds or resonance energies and their widths. Nevertheless, such methods have been used in the description of astrophysically relevant processes by allowing adjustments of the nucleon–nucleon interaction parameters, by renormalizing bound and resonance energies or astrophysical S–factors. Review articles on microscopic theories and/or their application to astrophysical processes are found in (3)–(7).

The potential model is based on the description of the dynamics of nuclear processes by a Schrödinger equation with local potentials in the entrance and exit channels and is appropriate for the description of direct reactions. Potential models can be applied to elastic scattering (Optical Model: OM), to transfer reactions (Distorted Wave Born Approximation: DWBA) and to capture reactions (Direct Capture: DC). The most important ingredients of potential models are the wave functions for the scattering, bound and quasi–bound (resonant) states in the entrance and exit channels. In the work performed by our group these wave functions have been determined from potentials calculated by the folding procedure (5), (8), (9). The nuclear densities for the folding

procedure are derived from nuclear charge distributions with an energy– and density–dependent NN–interaction. Important for the success of the potential models is the fact that the strength of the potentials are fitted to experimental elastic scattering cross sections and to the energies of bound and quasi–bound states. In this respect the potential models together with the folding procedure combine the first–principle approach of a microscopic theory with the flexibility of a phenomenological method.

NUCLEAR REACTIONS IN SOLAR HYDROGEN BURNING

In this section we want to discuss calculations of reactions taking place in solar hydrogen burning. We concentrate on reactions determining the branching ratios of the pp–chain, which are of interest for the solar neutrino problem. These reactions proceed predominately through the direct reaction mechanism. The main solar neutrino flux (95% for the ^{37}Cl experiment, 100% for the Kamiokande experiment and 92.5% for the ^{71}Ga experiments) come from the three pp–chains (Figs. 1 and 2). The nuclear reactions which determine the branching ratios of the pp–chain are the following (Fig. 1):

1. The reactions ^{3}He(^{3}He,2p)^{4}He and ^{3}He(α,γ)^{7}Be determine the branching ratio between the ppI– and (ppII+ppIII)–chains.

2. The reactions ^{7}Be(p,γ)^{8}B and ^{7}Be(e$^-$,ν)^{7}Li determine the branching ratio between the ppII– and ppIII–chains.

A general overview of the solar neutrino problem up to 1988 containing also the relevant nuclear processes can be found in the book of Bahcall (11). The main emphasis of this article is on recent theoretical calculations of astrophysical S–factors for the above reactions and their comparison with the currently accepted low–energy values (12) and recent experimental data.

The reactions ^{3}He(^{3}He,2p)^{4}He and ^{3}He(α,γ)^{7}Be

Since the reactions ^{3}He(α,γ)^{7}Be and ^{3}He(α,γ)^{7}Be are important for determining the branching ratio of the ppI– and (ppII+ppIII)–chains in solar hydrogen burning, the magnitudes of the cross sections for both reactions are of special interest for the solar neutrino problem (11).

The reaction ^{3}He(^{3}He,2p)^{4}He has been measured by a number of authors in the subCoulomb energy range (13)–(16), most recently by Krauss et et al. (17). In the later work the astrophysical S–factor was measured in the energy range $E_{\text{c.m.}} = 17.9$–342.9 keV. These data extend into the the thermonuclear energy region of the sun. In agreement with other experimental results ((18) and references therein) no evidence for a suggested low–energy resonance (19), (20)

FIGURE 1. The pp–chain (adapted from (10)).

FIGURE 2. The neutrino flux as a function of the neutrino energy from the different neutrino reactions in solar hydrogen burning. On the upper scale also the thresholds of the earth–bound neutrino experiments are shown (adapted from (10)).

FIGURE 3. Astrophysical S–factor for the reaction ^{3}He(^{3}He,2p)^{4}He calculated from a microscopic calculation (22) (short–dashed curve) and the potential model (23) (long–dashed curve) and compared to the currently accepted value low–energy values (12) (solid line) and experimental data [circles (15)], [triangles (17)].

has been found. Such a low–energy resonance would significantly enhance the ^{3}He+^{3}He–route in the pp–chain (21).

From the above measurements the data has been extrapolated to the thermonuclear energy range using either phenomenlogical models with quadratic polynomials (17), microscopic methods like the coupled–channel RGM (22) or the potential model (23) (Fig. 3).

In the RGM (22) as well as the DWBA (23) calculation the exit channel of the reaction ^{3}He(^{3}He,2p)^{4}He is treated by the formation of an α–particle and a diproton cluster. In both calculations parameters are adjusted to properties of the combined cluster sytem in such a way that it fits reasonably the low–energy experimental reaction data.

Experimental data for the ^{3}He(α,γ)^{7}Be cross section at subCoulomb energies has been obtained by numerous experimental groups (24)–(29). Capture cross sections observing the decay of ^{7}Be residual nucleus have been measured in further experiments (27), (29)–(31). The experimental data of Krähwinkel et al. (26) have been renormalized later on by a factor of 1.4 as suggested by Hilgemeier et al. (29).

Theoretically, Tombrello and Parker (32) have first succeeded in describing the energy dependence of the astrophysical S–factor in a direct–capture model using a hard–sphere potential. Further calculations in the framework of the potential model have been carried out in (33). The scattering wave functions of ^{3}He and ^{4}He and the bound–state wave functions of ^{7}Be were constructed by a phenomenological Woods–Saxon potential model and the orthogonality condition model. These models account for the measured elastic scattering and excitation energies of the low–lying states of ^{7}Be. In a further approach (34) for the lowest states in A=7 nuclei the model of a real ^{3}He nucleus interacting with a α–particle through a deep local potential of Gaussian shape is adopted. The four parameters of the model are determined by fitting them to reproduce numerous independent data in ^{7}Be.

Analyses using microscopic theories based on the RGM have been performed in (35)–(40). In the work of Kajino and arima(36) no normalization is used in order to obtain the astrophysical S–factor. The Pauli principle is fully taken into account. The radiative capture and the elctromagnetic properties using a multichannel RGM which also allows for studying the influence of inelastic scattering and distortion effects has been investigated in (37). For the relative cluster motion a number of Gaussian functions with different width parameters have been used. The capture reaction ^{3}He$(\alpha,\gamma)^7$Be was also studied in the framework of a microscopic potential model which is justified on the basis of microscopic many–body scattering theories (38).

In our calculation of the reactions ^{3}He$(^3$He,2p$)^4$He and ^{3}He$(\alpha,\gamma)^7$Be in the potential model the folding potentials are adjusted to reproduce the differential cross sections of the ^{3}He–^{3}He and ^{3}He–^{4}He elastic scattering measured in the range 1–3 MeV. The fit to the cross section data results in a strong parity dependance of the potentials. The ^{3}He–^{3}He potentials also describe very well the elastic scattering phase shift calculated in the RGM model ((23) and references therein) up to an energy of $E_{\mathrm{lab}} = 30$ MeV. In the ^{3}He–^{4}He case good agreement is found between the measured cross sections and the experimental and calculated elastic phase shifts up to $E_{\mathrm{lab}} \approx 12$ MeV ((41) and references therein).

The optical potentials obtained were used to calculate the astrophysical S–factors of the transfer reaction ^{3}He$(^3$He,2p$)^4$He and the capture reaction ^{3}He$(\alpha,\gamma)^7$Be. For the diproton system we choose a Fourier–Bessel charge–distribution. The strenght of the ^{4}He–2p potential was adjusted to reproduce the experimental reaction cross section at 150 keV. The spectroscopic factors for the above calculations were taken from shell–model calculations.

In Figs. 3 and 4 the results of our calculations are given by the long–dashed curves. For both reactions the agreement between the experimental and calculated data is excellent. The S_0–factors obtained in our calculations agree very well with the currently accepted values (12) if the experimental

FIGURE 4. Astrophysical S–factor for the reaction ^{3}He$(\alpha,\gamma)^7$Be calculated from different microscopic calculations (35)–(39) (upper and lower limits are shown by short–dashed curves) and the potential model (41) (long–dashed curve) and compared to the currently accepted low–energy values (12) (solid line) and experimental data [circles (26), triangles (27)].

errors are taken in consideration (Table 1).

Compared to the microscopic calculations the potential–model calculations are not dependent on different adopted model spaces and effective nucleon–nucleon interactions. Instead the relevant scattering and bound–state wave functions are determined uniquely from elastic scattering data and resonant and bound–state energies (except for the ^{3}He–2p system, where no experimental data is available). For the ^{3}He$(^3$He,2p$)^4$He reaction the microscopic calculation (22) underestimates the data at at higher energies, when compared to the experimental data (15), (17), to the potential model calculation (23) and to the currently accepted low–energy values (12) (Fig. 3). The reason for this deficiency is probably that in the microscopic calculation only s– and d–waves in the entrance channel are included, whereas in the potential model partial waves up to $\ell = 6$ are taken into account. For the ^{3}He$(\alpha,\gamma)^7$Be reaction the energy dependance of the different microscopic and the potential–model calculations are quite similar. However, for the absolute value the astrophysical S–factor for the different microscopic calculations scatter within about a factor of 2

Table 1. Recently calculated low–energy astrophysical S–factors for reactions of the pp–chain (upper part) compared to the currently accepted values of Bahcall (12) (lower part).

Reaction	Reference	S_0 [keV b]	dS/dE [b]
$^3\mathrm{He}(^3\mathrm{He,2p})^4\mathrm{He}$	(23)	$4.854 \cdot 10^3$	-1.328
$^3\mathrm{He}(\alpha,\gamma)^7\mathrm{Be}$	(41)	0.516	$-3.67 \cdot 10^{-4}$
$^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$	(47)	0.0249	$-3.2 \cdot 10^{-5}$
$^3\mathrm{He}(^3\mathrm{He,2p})^4\mathrm{He}$	(12)	$5.0(1 \pm 0.18) \cdot 10^3$	-0.9
$^3\mathrm{He}(\alpha,\gamma)^7\mathrm{Be}$	(12)	$0.533(1 \pm 0.096)$	$-3.1 \cdot 10^{-4}$
$^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$	(12)	$0.0224(1 \pm 0.28)$	$-3 \cdot 10^{-5}$

(Fig. 4). This is mainly due to the different adopted effective nucleon–nucleon interactions in the microscopic calculations. Again the potential–model gives an unique value, because the wave functions are determined from the elastic scattering data and resonant and bound–state energies. Contrary to the potential–model only the lower limit of the microscopic calculations can reproduce the experimental data (26), (27) and the currently accepted low–energy values (12) (Fig. 4).

The reactions $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$ and $^7\mathrm{Be}(\mathrm{e}^-,\nu)^7\mathrm{Li}$

The most uncertain of all nuclear reaction rates relevant for the solar neutrino problem is the reaction $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$. The reason is that there exist two measurements at low energies which differ by about 30% in the absolute magnitude for the astrophysical S–factor (42), (43). Both measurements find about the same low–energy dependance of the astrophysical S–factor in accordance with the microscopic calculations (44)–(46) and the potential–model approach (47). However, due to different adopted model spaces and effective nucleon–nucleon interactions also the absolute values of the astrophysical S–factors for the microscopic calculations differ considerably (Fig. 5). Therefore, the absolute values of the astrophysical factors obtained from the microscopic calculations have either been fitted to the two existing experimental data sets (45) or a range of different possible values for the astrophysical S–factor is presented by the authors (46).

In the potential–model approach for the reaction $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$ the strength of the folding potentials are adjusted to reproduce the resonant and bound states in the entrance and exit channel (47). This ensures that the asymptotic parts of the wave functions are correct. This is of special importance, because the main contributions for this reaction come from far out in the nuclear exterior (for the reaction $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$ at 15 keV the main contributions come from a region of about 40 fm outside the target nucleus) (47). The energy dependence of

FIGURE 5. Astrophysical S–factor for the reaction ^{7}Be(p,γ)^{8}B calculated from different microscopic calculations (44)–(46) (upper and lower limits are shown by short–dashed curves) and the potential model (47) (long–dashed curve) and compared to the currently accepted low–energy values (12) (solid line) and experimental data [squares (42), diamonds (43)].

the astrophysical factors is consistent with the experimental data sets as well as the microscopic calculations. However, using the wave functions obtained with the folding procedure it is also possible to calculate the absolute values of the astrophysical S–factor for this reaction without using a renormalization procedure (47). This calculation reproduces almost the data of Kavanagh et al. (42) (Fig. 5) and therefore favours the higher experimental value of the astrophysical S–factor. The potential–model calculation is about 11% higher than the currently accepted value (12) (Table 1 and Fig. 5).

Recent progress has also been made on the electron capture rate by a ^{7}Be–nucleus through the reaction ^{7}Be(e$^-$,ν)^{7}Li in the solar plasma by performing a self–consistent study of continuum and bound electrons (45). The improved treatment of the screening effects results in a small increase for this reaction rate of about 1.3% compared to the currently accepted value given in (11).

Table 2. Observed solar neutrino rates compared to the predictions of the SSM. (From (12), (48)).

Experiment	Observation (SNU)[a]	Prediction (SNU)[a]	$\frac{\text{Observation}}{\text{Prediction}}$ [b]
Davis	2.2 ± 0.2	8.0 ± 3.0	0.275 ± 0.025
Kamiokande			$0.47 \pm 0.05 \pm 0.06$
Gallex	$83 \pm 19 \pm 8$	131.5^{+20}_{-17}	0.63 ± 0.20
Sage	$58^{+17}_{-24} \pm 14$	131.5^{+21}_{-17}	$0.44^{+0.24}_{-0.29}$

[a] 1 SNU = 1 Solar Neutrino Unit = 10^{-36} ν–captures per s and target atom.

[b] The errors reflect only the observational errors.

NUCLEAR ASPECTS OF THE SOLAR NEUTRINO PROBLEM

The neutrino–producing nuclear reactions in the sun are shown in Fig. 2. As can be seen from this figure the strongest neutrino flux stems from the reaction $p + p \rightarrow d + e^+ + \nu$. However, these neutrinos have very low energies $E_\nu \leq 0.42\,\text{MeV}$. On the other hand almost all the neutrinos with $E_\nu > 2\,\text{MeV}$ are produced in the decay of ^{8}Be. The neutrino energies play a crucial role in their detection, because all neutrino detectors have certain thresholds (see upper scale in Fig. 2). The threshold energies are $0.814\,\text{MeV}$ for the Davis (^{37}Cl–detector), about $7.5\,\text{MeV}$ for the Kamiokande (neutrino–electron scattering) and $0.233\,\text{MeV}$ for the Gallex and Sage experiments (^{71}Ga–detector).

The predicted neutrino flux of the SSM is compared to the different observed solar neutrino rates from the Davis, Kamiokande, Gallex and Sage experiments in Table 2. Roughly speaking one can say that the observed neutrino flux is about 1/3, 1/2 and 2/3 of the predicted neutrino flux for the Davis, Kamiokande, and Gallex and Sage experiments, respectively. In the following we want to investigate the dependance of the predicted fluxes for the different solar neutrino experiments on the nuclear input using a simple model. The fluxes ϕ for the different solar neutrino sources of the pp–chain can be expressed through the following proportionality relations (assuming constant luminosity, radius, metal abundance and age of the sun) (11):

$$\phi(\text{pp}) \propto \phi(\text{pep}) \propto S_{11}^{0.14} S_{33}^{0.03} S_{34}^{-0.06} \tag{1}$$

$$\phi(^7\text{Be}) \propto S_{11}^{-0.97} S_{33}^{-0.43} S_{34}^{0.14} \tag{2}$$

$$\phi(^8\text{B}) \propto S_{11}^{-2.6} S_{33}^{-0.40} S_{34}^{0.81} S_{17}^{1.0} \tau_{e7}^{1.0} \tag{3}$$

In these equations the S denote the low–energy astrophysical S–factors for the different reactions wich are characterized by:

pp: $p + p \rightarrow d + e^+ + \nu$

pep: $p + e^- + p \rightarrow d + \nu$

33: $^3\mathrm{He}(^3\mathrm{He},2\mathrm{p})^4\mathrm{He}$
34: $^3\mathrm{He}(\alpha,\gamma)^7\mathrm{Be}$
17: $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$
e7: $^7\mathrm{Be}(\mathrm{e}^-,\nu)^7\mathrm{Li}$.

Compared to the formulae given by Bahcall (11) we multiplied the reaction rate for the $^8\mathrm{B}$–neutrino flux additionally with the factor τ_{e7}. This factor describes the lifetime of $^7\mathrm{Be}$ with respect to the reaction $^7\mathrm{Be}(\mathrm{e}^-,\nu)^7\mathrm{Li}$. Since the concentration of $^7\mathrm{Be}$ is proportional to its lifetime with respect to electron capture of $^8\mathrm{Be}$ (the lifetime with respect to $^7\mathrm{Be}(\mathrm{p},\gamma)^8\mathrm{B}$ is much longer) the $^8\mathrm{B}$–neutrino flux is proportional to the lifetime τ_{e7} as well as to S_{17}.

For the different neutrino–experiments we can deduce the following expressions for the predicted solar neutrino fluxes in the SSM for the different relevant reaction rates of the pp–chain:

1. $^{37}\mathrm{Cl}$ experiment using (12):

$$\Phi(^{37}\mathrm{Cl}) \;=\; (0.2R(\mathrm{pep}) + 1.2R(^7\mathrm{Be}) + \tag{4}$$
$$+6.2R(^8\mathrm{B}) + 0.4)\,\mathrm{SNU}$$

2. Kamiokande experiment:

$$\frac{\Phi(\mathrm{Kamiokande(Observation)})}{\Phi(\mathrm{Kamiokande(Prediction)})} = 0.47R(^8\mathrm{B}) \tag{5}$$

3. $^{71}\mathrm{Ga}$ experiments:

$$\Phi(^{71}\mathrm{Ga}) \;=\; (70.8R(\mathrm{pp}) + 3.1R(\mathrm{pep}) + 35.8R(^7\mathrm{Be}) + \tag{6}$$
$$+13.8R(^8\mathrm{B}) + 7.9)\,\mathrm{SNU}$$

In the above equations the quantities R denote the reaction rates (or low–energy astrophysical S–factors) of eqs. (1)–(3) normalized to the currently accepted values given in (12). We want to emphasize that eqs. (1)–(6) can only give a semiquantitative understanding of the dependence of the solar neutrino fluxes, because the actual variations of the calculated neutrino fluxes are determined by the coupled partial differential equations of stellar evolution and the boundary conditions (48). However, in many cases the formulae (1)–(6) can give a first approximation of the dependence of the solar neutrino fluxes on the astrophysical S–factors.

Using the new input rates discussed in the foregoing section we can use the above formula to recalculate the neutrino input fluxes from the individual neutrino fluxes as well as the changes in the predicted rate for the neutrino experiments. The pp neutrino flux is almost unchanged, the $^7\mathrm{Be}$ and $^8\mathrm{B}$ neutrino fluxes increase by 1% and 11%, respectively, when compared them to

the currently accepted values (12). The enhancement of the ^{8}B neutrino flux is due to the higher astrophysical S–factor for the reaction ^{7}Be(p,γ)^{8}B which is the most uncertain nuclear reaction rate with respect to the solar neutrino problem. With the new nuclear reaction rates we get the following predictions for the neutrino fluxes enforcing the solar neutrino problem to some extend:

1. ^{37}Cl experiment: $\Phi(^{37}\text{Cl}) = 8.7\,\text{SNU}$
   (instead of 8.0 SNU (12)).

2. Kamiokande experiment: $\frac{\Phi(\text{Kamiokande(Observation)})}{\Phi(\text{Kamiokande(Prediction)})} = 0.42$
   (instead of 0.47 (12)).

3. ^{71}Ga experiments: $\Phi(^{71}\text{Ga}) = 133.6\,\text{SNU}$
   (instead of 131.5 SNU (12)).

We can also use the above formulae to use an astrophysical S–factor for the reaction ^{7}Be(p,γ)^{8}B determined from the recent experiment at RIKEN measuring the low–energy ^{8}B + ^{208}Pb $\rightarrow$ p + ^{7}Be + ^{208}Pb Coulomb dissociation cross section (48). Taking the E1 and E2 contributions into account the astrophysical S–factor for the reaction ^{7}Be(p,γ)^{8}B is only 55% (49) of the current accepted value (12). In this case we obtain

1. ^{37}Cl experiment: $\Phi(^{37}\text{Cl}) = 5.2\,\text{SNU}$
   (instead of 8.0 SNU (12)).

2. Kamiokande experiment: $\frac{\Phi(\text{Kamiokande(Observation)})}{\Phi(\text{Kamiokande(Prediction)})} = 0.85$
   (instead of 0.47 (12)).

3. ^{71}Ga experiments: $\Phi(^{71}\text{Ga}) = 123.9\,\text{SNU}$
   (instead of 131.5 SNU (12)).

This brings the predicted neutrino flux of the high–energy neutrinos within the experimental errors of the Kamiokande experiment shown in Table 1. However, the predicted neutrino flux of the ^{37}Cl experiment of 5.2 SNU is still outside the experimental errors of this experiment. This result is also approximately obtained by a much more sophisticated method using 1000 solar models with modified fluxes in a Monte Carlo simulation (50). None of these considered models lie within 3σ measurement errors of the ^{37}Cl–experiment. This means that at the present time it is still not possible to reconcile the ^{37}Cl and Kamiokande experiments within the framework of the SSM.

Electron screening plays an important role in low–energy laboratory measurements of nuclear reactions as well as in the solar plasma. However, the physics behind the screening effects for these two cases is quite different. In laboratory experiments the bound electrons are inevitably present in the target

and shield part of the barrier of the bare nuclei, whereas in the plasma the electrons occupy continuum states. At present there is a significant, unexplained discrepancy between the shielding effect in laboratory experiments and the theoretical predictions for the screening energies (the effective energy increase through screening) which are to low up to a factor of about 2 when compared to experiment (7), (51). However, the knowledge of screening effects is absolutely necessary in order to make reliable extrapolations of experimentally determined astrophysical S–factors to the thermonuclear energies of the sun. The above discussed discrepancy which may also be due to some unknown nuclear effect raises doubts about the correctness of these extrapolations to thermonuclear energies.

SUMMARY

The solar neutrino problem has persisted pertinaciously solutions within standard physics. In this respect the nuclear reaction rates obtained from laboratory efforts by measuring solar nuclear reactions in earth–bound accelerator experiments or theoretical analyses are no exceptions. The calculated reaction rates which are discussed in this work even enforces the solar neutrino problem to some extend, because they favour the higher astrophysical S–factor for the reaction $^7\text{Be}(\text{p},\gamma)^8\text{B}$. Even if the reaction rate for this reaction is changed drastically as measured by the recent experiment at RIKEN, one can reproduce the Kamiokande experiment within the experimental errors, but it seems not to be possible to obtain consistency with the ^{37}Cl experiment. At the present time there exists stil a great uncertainty in the extrapolation of the astrophysical S–factors to thermonuclear energies due to the incorrect treatment of screening effects or some other unknown nuclear effect.

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